

Forward-backward $t\bar{t}$ asymmetry from anomalous stop pair production

Gino Isidori^{1,*} and Jernej F. Kamenik^{2,3,†}

¹*INFN, Laboratori Nazionali di Frascati, Via E. Fermi 40 I-00044 Frascati, Italy.*

²*J. Stefan Institute, Jamova 39, P. O. Box 3000, 1001 Ljubljana, Slovenia*

³*Department of Physics, University of Ljubljana, Jadranska 19, 1000 Ljubljana, Slovenia*
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We analyse a simple Standard Model (SM) extension with only two new light fields: a scalar partner of the top \tilde{t} (with mass above m_t) and a light neutral fermion χ^0 (with mass of a few GeV), coupled to SM quarks via a Yukawa interaction. We show that such model can lead to a significant enhancement of the forward-backward asymmetry in $t\bar{t}$ production at Tevatron via the additional $t\bar{t}$ pairs produced from $\tilde{t}\tilde{t}^\dagger$ decays. The model satisfies existing constraints on new-physics searches both at low and high energies, and could even address the cosmological dark-matter abundance. The implications for future searches at the LHC are briefly outlined.

I. INTRODUCTION

There are recent experimental indications of an anomalously large forward-backward asymmetry (FBA) in top-antitop pair production at the Tevatron. The asymmetry is defined as

$$A_{FB}^{t\bar{t}} = \frac{\sigma_{\Delta y > 0} - \sigma_{\Delta y < 0}}{\sigma_{\Delta y > 0} + \sigma_{\Delta y < 0}}, \quad (1)$$

where $\Delta y = y_t - y_{\bar{t}}$ is the difference in rapidity, and positive rapidity is measured in the direction of the colliding p . The most significant measurement of $A_{FB}^{t\bar{t}}$ is due to the CDF collaboration which after subtracting backgrounds and performing the unfolding procedure to recover the asymmetry at the partonic level reports [1]

$$A_{FB}^{t\bar{t}} = 0.158 \pm 0.074, \quad (2)$$

to be compared with the NLO QCD prediction $(A_{FB}^{t\bar{t}})_{SM} = 0.058 \pm 0.009$ [2]. The observed FBA is even larger at high invariant masses of the $t\bar{t}$ system, in particular

$$A_h = A_{FB(m_{t\bar{t}} > 450 \text{ GeV})}^{t\bar{t}} = 0.475 \pm 0.114, \quad (3)$$

while QCD predicts 0.088 ± 0.013 [2].

At the same time, the total inclusive top pair production cross-section is currently measured to be [3]

$$\sigma^{t\bar{t}} = (7.50 \pm 0.48) \text{ pb}, \quad (4)$$

that is consistent with the most recent theoretical SM predictions for this observable: $(7.2 \pm 0.4) \text{ pb}$ [4] and $(6.4 \pm 0.4) \text{ pb}$ [5]. Another important experimental constraint is the $m_{t\bar{t}}$ distribution of the production cross-section. As pointed out in [6], the most significant information at high $m_{t\bar{t}}$ is the one derived from the next-to-highest measured bin [7]

$$\sigma_h = \sigma_{(700 \text{ GeV} < m_{t\bar{t}} < 800 \text{ GeV})}^{t\bar{t}} = (80 \pm 37) \text{ fb}, \quad (5)$$

to be compared with $(\sigma_h)_{SM} = 80 \pm 8 \text{ fb}$ [5].

Several authors have analysed the possibility that the large FBA reported by CDF is obtained by an anomalous $t\bar{t}$ production mechanism interfering with the corresponding SM process [6, 8]. The interference of the two amplitudes maximizes the possible impact on $A_{FB}^{t\bar{t}}$ while minimizing deviations from the SM in $\sigma^{t\bar{t}}$. However, we observe that the present cross-section measurements still leave some room for an additional (incoherent) production of $t\bar{t}$ pairs. This motivates us to analyse a different mechanism to enhance $A_{FB}^{t\bar{t}}$, namely an anomalous production of $t\bar{t}$ + invisible particles, which do not interfere with any SM process.

At the level of the inclusive measurement, the non-standard production (passing the experimental selection cuts) can still contribute of up to 13% compared to the SM (using the higher SM prediction in [4] as a conservative reference value). The percentage of the new contribution could even rise at high $m_{t\bar{t}}$, provided it does not exceed 50% of the SM cross-section in σ_h . Ideally, if the anomalous production would carry a 100% FBA, it could perfectly accommodate the measured value of the total FBA without violating the bound on the cross section.

The production of the $t\bar{t}$ + invisible final state can be obtained by the pair production of “top partners” – particles decaying into a top quark and light invisible states – that are naturally expected in several SM extensions. If the mass difference between the top partner and the top is sufficiently small, the missing energy carried by the invisible states is small and the $t\bar{t}$ pairs thus produced would pass the experimental cuts applied to identify $t\bar{t}$ pairs in the SM. A prototype of such scenario would be the fourth generation up quark or a vector-like fermionic top partner, decaying into a top and one or several light invisible particles. However, colored fermions have large QCD cross-sections with a vanishing FBA contribution at leading order: this makes them unattractive candidates for our purpose. On the contrary, a scalar top partner (\tilde{t}) of mass around 200 GeV decaying into a top and a single invisible particle (χ^0) is still perfectly allowed [9]. This is because the QCD production cross-section for scalar particles proceeds mainly via p -wave

* Electronic address: gino.isidori@lnf.infn.it

† Electronic address: jernej.kamenik@ijs.si

and thus vanishes at threshold [10].

The stop and a light neutralino are naturally present in supersymmetric extensions of the SM, making this scenario particularly attractive. As we will discuss below, in the minimal supersymmetric extension of the SM (MSSM) the FBA generated by this mechanism is vanishingly small, but it could become sizable in more general frameworks. Without assuming a specific model, here we adopt a phenomenological bottom-up approach: we assume \tilde{t} and χ^0 to be the only light relevant non-standard particles, and we determine the nature of their interactions mainly by looking at phenomenological constraints. As a remnant of R -parity in the MSSM, we also assume a discrete Z_2 symmetry under which only these non-standard states are charged, such that they can only be produced in pairs and such that χ^0 is a stable particle.

II. IDENTIFICATION OF THE MODEL

In order to generate a large $A_{FB}^{t\bar{t}}$, the differential partonic cross-section $u(p_1)\bar{u}(p_2) \rightarrow \tilde{t}(p'_1)\tilde{t}^\dagger(p'_2)$ should exhibit a large \hat{t} -odd dependence,¹ where $\hat{t} = (p'_1 - p_1)^2$. This dependence is not generated by the leading QCD contribution to $u\bar{u} \rightarrow \tilde{t}\tilde{t}^\dagger$.

A first approach to generate a sizable \hat{t} -odd dependence in $u\bar{u} \rightarrow \tilde{t}\tilde{t}^\dagger$ is via heavy mediators (in either \hat{s} or \hat{t} channel). Integrating out the heavy mediators leads to an appropriate set of higher-dimensional effective operators coupling the up quarks to the $\tilde{t}\tilde{t}^\dagger$ pair. Up to canonical dimension six, the relevant operators are

$$\bar{u}u\tilde{t}^\dagger\tilde{t}, \quad \bar{u}\gamma_5 u\tilde{t}^\dagger\tilde{t}, \quad \bar{u}\gamma_\mu u\tilde{t}^\dagger\partial^\mu\tilde{t}, \quad \bar{u}\gamma_\mu\gamma_5 u\tilde{t}^\dagger\partial^\mu\tilde{t}. \quad (6)$$

In this case the required \hat{t} -odd dependence can only appear in the numerator of the cross section, but the kinematical structure of all the above operators is such that they do not generate it. A non-vanishing \hat{t} dependence can appear only introducing operators of dimension 7 or higher, which are expected to be more suppressed by naive dimensional analysis. Since the contributions with a \hat{t} -odd dependence should dominate the cross section, which receives also a sizable contribution from ordinary QCD, this approach appears to be highly contrived.

A second approach is to assume light \hat{t} -channel mediators for $u\bar{u} \rightarrow \tilde{t}\tilde{t}^\dagger$. In this way the \hat{t} -odd dependence is naturally induced by the \hat{t} -channel propagators and gets more pronounced the lighter the mediators. The lightness of the mediators is also required to generate a sizable cross-section (larger than the QCD induced one) while maintaining perturbativity of the associated mediator couplings to (light) quarks and \tilde{t} . The mediators need

to be neutral and, more generally, $SU(2)_L \times U(1)_Y$ singlets, to avoid the bounds on new light states from LEP. Assuming they couple trilinearly to u and \tilde{t} , they should also be fermions. Since we have already introduced a light neutral fermion in order to account for $\tilde{t} \rightarrow t\chi^0$, the simplest possibility is to assume χ^0 itself to be the mediator. We are thus led to consider the following simple Lagrangian

$$\mathcal{L} = \mathcal{L}_{SM} + (D_\mu\tilde{t})^\dagger(D^\mu\tilde{t}) - m_{\tilde{t}}^2\tilde{t}^\dagger\tilde{t} + \bar{\chi}^0(i\gamma_\mu D^\mu)\chi^0 - m_\chi\bar{\chi}_c^0\chi^0 + \sum_{q=u,c,t} (\tilde{Y}_q\bar{q}_R\tilde{t}\chi^0 + \text{h.c.}), \quad (7)$$

where we have introduced effective $q_R\tilde{t}\chi^0$ couplings for all three generations of right-handed up quarks. Having assumed χ^0 to be an $SU(2)_L \times U(1)_Y$ singlet, this implies that \tilde{t} has the same $SU(2)_L \times U(1)_Y$ quantum numbers of right-handed up quarks. As we will discuss below, beside the minimality of the new states introduced, the choice of a pure right-handed coupling minimizes the impact on electroweak observables and flavour-changing neutral-current (FCNC) transitions.

Following the principle of introducing the minimal set of relevant new states, we have assumed χ^0 to have only a Majorana mass term. This hypothesis has essentially no impact on collider phenomenology (provided m_χ is sufficiently small in order to suppress the production of same-sign tops), while it may be an important ingredient if we require χ^0 to be a viable dark-matter candidate (see below).

With this choice \tilde{t} can be identified with the right-handed stop of the MSSM or, more generally, with any combination of right-handed up-type squarks. In principle, χ^0 could be identified with the bino of the MSSM. However, bino couplings are completely determined by electroweak gauge symmetries and ultimately turn out to be too small to significantly affect the FBA. On the other hand, we see no obstacles to accomodate the required framework in extensions of the MSSM with extra $SU(2)_L \times U(1)_Y$ singlets. An important condition in order to implement this mechanism in supersymmetric extensions of the SM is a sufficiently heavy gluino mass, such that gluino-mediated amplitudes do not provide a significant enhancement of the $\tilde{t}\tilde{t}^\dagger$ production cross section.

As anticipated, we keep our discussion general analysing the phenomenological implications of (7) without assuming extra model-dependent conditions. In order to avoid the production of χ^0 pairs in charm decays we require $m_\chi > m_D/2$, while for the \tilde{t} state we require $m_{\tilde{t}} > m_t$ and we impose the constraints derived in [9] for top-partner pair production at Tevatron.

III. COLLIDER PHENOMENOLOGY

As an illustration of the resulting phenomenology at the Tevatron, we choose $m_{\tilde{t}} = 200$ GeV and $m_{\chi^0} =$

¹ The $u\bar{u}$ initial partonic state is chosen because it largely dominates over $d\bar{d}$ at the Tevatron, especially at high invariant masses, while the $g\bar{g}$ initial state cannot produce a FBA.

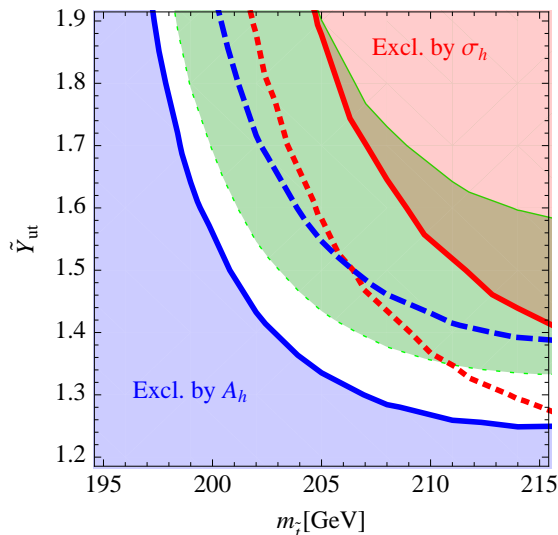


Figure 1. Tevatron constraints in the $m_{\tilde{t}}-\tilde{Y}_u$ plane. The inclusive $A_{FB}^{t\bar{t}}$ and $\sigma^{t\bar{t}}$ are reproduced within 1σ in the central green band. The region below the continuous (dashed) blue line is excluded by A_h at 95% C.L. (90% C.L.). The region above the continuous (dotted) red line is excluded by σ_h at 95% C.L. (90% C.L.).

2 GeV. We simulate our signal $p\bar{p} \rightarrow t\bar{t} \rightarrow t\bar{t}\chi^0\chi^0$ at the partonic level with MadGraph/MadEvent 4.4.57 [11] and using CTEQ6L1 set of PDFs [12]. For $\tilde{Y}_u \simeq 0.85$ we find that the total $t\bar{t}$ production cross-section reaches 12% of the SM $t\bar{t}$ cross-section, exhibiting a 50% FBA. Assuming \tilde{t} decays dominantly into $t\chi^0$ (or assuming $\tilde{Y}_t \gg \tilde{Y}_{u,c}$) and applying the $t\bar{t}$ reconstruction cuts, this leads to a predicted $A_{FB}^{t\bar{t}}$ within one standard deviation from the experimental data in (2). At the same time, with this parameter choice the value of the total cross-section is within one standard deviation from the measured one, using the prediction in [4] as a conservative SM normalization. Some tension does develop only when comparing the high invariant-mass data, where the predicted cross-section and FBA tend to be higher and smaller than data, respectively. In particular, the predicted σ_h exhibits a tension of 1.6σ when compared with (5), while the SM tension of 3.4σ in A_h is only relieved to around 1.9σ .

In general, as long as χ^0 is much lighter than \tilde{t} and $\tilde{Y}_t \gg \tilde{Y}_{u,c}$, the results do not depend on the precise values of m_{χ^0} and \tilde{Y}_t . For this reason, we have performed a more accurate evaluation of the various constraints on the parameter space of the model at fixed m_{χ^0} and \tilde{Y}_t . In particular, we have set $m_{\chi^0} = 2$ GeV and $\tilde{Y}_t = 4$ (the largest value to ensure a perturbative behavior in \tilde{Y}_t) and varied the remaining relevant parameters, namely \tilde{Y}_u and $m_{\tilde{t}}$. Since the condition $\tilde{Y}_t \gg \tilde{Y}_u$ is not always fulfilled, we have taken into account the actual $\tilde{t} \rightarrow t\chi^0$ branching ratio when computing our signals. The resulting comparison with the relevant constraints coming from $t\bar{t}$ production at the Tevatron is displayed in Fig. 1.

As shown in Fig. 1, the requirements of a large FBA and a small impact on the cross section are partly in conflict, especially at high invariant masses. However, there is a region of the parameter space where all constraints are satisfied at the 90% C.L. This happens for $200 \lesssim m_{\tilde{t}} \lesssim 205$ GeV and $\tilde{Y}_u \gtrsim 1.5$ (part of the green area where the dotted line is above the dashed one). The reason for that lies in the interplay of the \tilde{Y}_u dependence in the χ^0 -mediated cross-section generating a large FBA and in $Br(\tilde{t} \rightarrow t\chi)$. In particular, for $|\tilde{Y}_u| \lesssim |\tilde{Y}_t|$ the \hat{t} -even QCD contribution is suppressed by $Br(\tilde{t} \rightarrow t\chi)$ more than the \hat{t} -odd contribution, which enables us to obtain a larger FBA value for a given contribution to the cross-section.

The two additional invisible particles in the final state result in extra missing energy compared to the SM $t\bar{t}$ production. This might result in (1) observable modifications of the various kinematical distributions used in the experimental analyses of the FBA (as well as the $t\bar{t}$ production cross-section and the top quark mass) to discriminate signal from background; (2) tensions between the leptonic, semileptonic and all-hadronic modes in the precise measurements of the $t\bar{t}$ cross-section and the top quark mass. Unfortunately, most present top quark analyses employ multivariate techniques which are difficult to reproduce without the full detector simulation. Nonetheless, we have estimated the size of such effects by comparing the leptons and jets p_T distributions as well as distributions of missing transverse energy (E_T^{miss}) and the scalar sum of transverse particle energies ($H_T = \sum_i E_T^i$) in the process $p\bar{p} \rightarrow t\bar{t}\chi^0\chi^0$ (at preferred values of the model parameters) with those of the SM $t\bar{t}$ production. We have considered these distributions separately for the leptonic mode with kinematical cuts used in [13], the semileptonic reconstruction channel described in [14] as well as the all-hadronic mode used recently for the top mass measurement in [15]. We have simulated both signals using MadGraph, interfaced with Pythia 6.4.14 [22] for showering and hadronization, and the CDF PGS for detector simulation with a cone ($dR = 0.4$) jet reconstruction algorithm. In the leptonic mode, we first observe that the fraction of the $t\bar{t}\chi^0\chi^0$ signal passing the kinematical cuts is comparable to the SM $t\bar{t}$ signal (the same is true also for the semileptonic mode). Regarding the kinematical distributions outlined above, the effects of the extra χ^0 's can be readily described as a shift in the distributions by few GeV. Since such shifts only affect the $O(10\%)$ new contribution to the total $t\bar{t}$ sample, the resulting effects turn out to be much smaller than the reported systematic and statistical uncertainties in these distributions. As a consequence, in di-leptonic sample we expect the modifications in the total measured $t\bar{t}$ cross-section or the top-quark mass to be smaller than the present statistical uncertainties. In the semileptonic case, the shift in the distributions is a bit more pronounced, of the order of 10 GeV, but the resulting change in the combined $t\bar{t}\chi^0\chi^0 + t\bar{t}$ signal sample distributions is still smaller

than the associated statistical uncertainties. We have also compared the reconstructed transverse leptonic W masses ($m_T^W = \sqrt{2E_T^{\text{miss}}E_T^\ell[1 - \cos(\phi_\ell - \phi_{E_T^{\text{miss}}})]}$, where ϕ_i are the corresponding azimuthal angles) and found good agreement between the distributions of the two signals. Most important, the distribution peak in the $t\bar{t}\chi^0\chi^0$ sample does not appear shifted with respect to the SM $t\bar{t}$ case. Therefore, we expect that resulting effects on $t\bar{t}$ cross-section or top-quark mass measurements using the semileptonic mode to be within the present statistical and systematic uncertainties. Finally, for the all-hadronic mode, we find that the requirement of no significant E_T^{miss} reduces the fraction of $t\bar{t}\chi^0\chi^0$ signal events passing the cuts to one third, compared to the SM $t\bar{t}$ sample (assuming an uncertainty in E_T^{miss} reconstruction of $\sigma_{E_T^{\text{miss}}} \simeq 0.5\sqrt{H_T}\text{ GeV}^{1/2}$ [16]). The resulting discrepancy in the measured total $t\bar{t}$ production cross-section compared to the semileptonic and leptonic modes for preferred values of our model parameters would be of the order of 7%, that is less than the present uncertainties of the individual analyses.

Recently top-pair production has also been measured at the LHC. The value reported in [17] for the total $t\bar{t}$ cross section in pp collisions at $\sqrt{s} = 7$ TeV center-of-mass energy is

$$\sigma_{LHC}^{t\bar{t}} = (194 \pm 78) \text{ pb}, \quad (8)$$

in agreement with the SM prediction of 158 ± 24 pb [18]. For our preferred choice of parameters the corresponding $t\bar{t}^\dagger$ production cross-section is only around 10 pb [10]², still below the present sensitivity of the LHC experiments.

The $t\bar{t}^\dagger$ pairs decaying to two up quarks and two χ^0 's contribute to the jets plus missing transverse energy (E_T^{miss}) signatures, searched for both at the Tevatron [19] and at the LHC [20, 21]. At present the most sensitive search is the ATLAS analysis of two jets plus E_T^{miss} [21], based on 35 pb^{-1} of data, and in particular the scenario A of [21], which is optimized for low-mass squark pair production. Here two jets with $p_T > 140, 40$ GeV respectively, and $E_T^{\text{miss}} > 100$ GeV are required. Additional kinematic cuts are also imposed, among which the most relevant are $m_{eff} = \sum_{i=1}^2 |p_T^i| + E_T^{\text{miss}} > 500$ GeV and $E_T^{\text{miss}}/m_{eff} > 0.3$. Under these conditions 87 events are found passing all the cuts, with a SM background estimation of 118 ± 42 . Again we have simulated our signal contribution using MadGraph, interfaced with Pythia 6.4.14 [22] for showering and hadronization and the ATLAS PGS for detector simulation with a k_T jet reconstruction algorithm. For the most interesting region of the model's parameter space, we find that the signal

cross-section passing the p_T and E_T^{miss} cuts alone is at the level of 1 pb or less, still below the present experimental sensitivity.

In principle, also the Tevatron analyses for single top production could be used to set bounds on the parameter space of our model. However, in this case it is very difficult to assess the sensitivity. Present experimental analyses employ sophisticated multivariate discriminants including tight constraints on the missing transverse energy distribution, which is required to agree with expectations from a single, leptonically decaying W boson in the final state [23]. Here we simply note that the $t\bar{t}^\dagger$ production is small compared to the main SM backgrounds, namely W +jets and $t\bar{t}$ production, which have a similar signature. Thus we do not expect very stringent constraints from the present single-top production searches.

If χ is a Majorana particle, then in principle one can expect a non-vanishing production of same-sign tops from the $uu \rightarrow t\bar{t}$ process; however, this amplitude is strongly suppressed by the smallness of m_χ . For $m_\chi = 2$ GeV and $m_{\tilde{t}} > m_t$ the production of same sign tops is suppressed by $(m_\chi/m_{\tilde{t}})^2 \sim 10^{-4}$ with respect to the χ -mediated stop anti-stop production (that is a small fraction of the SM $t\bar{t}$ cross section). Beside the $(m_\chi/m_{\tilde{t}})^2$ power suppression, the same-sign top cross section receives logarithmic enhancements both near threshold and at large dilepton invariant masses; however, these are compensated by the small probability distribution for the initial uu state. As a confirmation of this qualitative evaluation, we have performed a quantitative study of the same-sign top cross section for our preferred range of parameters finding $\sigma(tt)_{\text{Tevatron}} \sim 0.01$ fb and $\sigma(tt)_{\text{LHC } 7 \text{ TeV}} \sim 3$ fb: a level well below the present sensitivity.

IV. FLAVOUR PHYSICS AND DARK MATTER

A very significant constraint on \tilde{Y}_q comes from the $D^0 - \bar{D}^0$ mixing amplitude, to which χ^0 and \tilde{t} can contribute at one loop. Introducing the following low-energy effective dimension-six Hamiltonian

$$\mathcal{H}_{\text{eff}} = C_{ud}^R (\bar{c}_R \gamma_\mu u_R)^2, \quad (9)$$

the one-loop induced contribution by our Lagrangian in (7) is

$$C_{ud}^R = -\frac{1}{32\pi^2 m_{\tilde{t}}^2} (\tilde{Y}_c \tilde{Y}_u^*)^2. \quad (10)$$

Using the bound $|C_{ud}^R| < 1.2 \times 10^{-3} \text{ TeV}^{-2}$ [24], this implies

$$|\tilde{Y}_c / \tilde{Y}_u| < 0.06, \quad (11)$$

for our preferred values of $m_{\tilde{t}}$ and \tilde{Y}_u . This is certainly a fine-tuned condition, although it is comparable to the

² At the LHC, the QCD production through gg initial state completely dominates over the anomalous \tilde{t} -channel contributions for our choices of parameters.

hierarchies exhibited by the Yukawa couplings within the SM.

Another interesting consequence of our scenario are FCNC top quark decays to a light quark jet and missing energy (i.e. $t \rightarrow u\chi^0\bar{\chi}^0$). Neglecting the χ^0 mass dependence, the decay rate is given by

$$\Gamma(t \rightarrow u\chi^0\bar{\chi}^0) = \frac{|\tilde{Y}_t\tilde{Y}_u|^2 m_t^5}{6144\pi^3 m_{\tilde{t}}^4}. \quad (12)$$

For our illustrative choice of parameters, the branching ratio might reach the level of 10^{-3} , which could be within the projected LHC sensitivity [25].

Finally, it is interesting to note that stable fermions like χ^0 with mass of a few GeV, annihilating to light quark pairs via effective dimension six operators have been previously considered as dark-matter candidates [26]. In our case, assuming for the moment that χ^0 is a Dirac fermion, the low-energy coupling of χ^0 to light quarks induced by the \tilde{t} exchange leads to

$$\mathcal{L}_{\text{annih.}}^{\text{eff}} = \frac{|\tilde{Y}_u|^2}{4m_{\tilde{t}}^2} \bar{u}_R \gamma_\mu u_R \bar{\chi}^0 \gamma^\mu (1 - \gamma_5) \chi^0. \quad (13)$$

In this case the dominant contribution to the thermal annihilation rate of χ^0 comes from the vector current part ($\bar{\chi}^0 \gamma_\mu \chi^0$) of the above operator. Using the results of [26], and setting $m_{\tilde{t}} \simeq 200$ GeV and $|\tilde{Y}_u| \simeq 1.5$ in (13) we find that the correct relic abundance of χ^0 is reproduced for $m_{\chi^0} \simeq 2$ GeV, that would perfectly fit with the collider data. Such contributions are however in some tension with existing direct dark-matter detection experiments [27]. Furthermore, a recent dark-matter search using the WMAP CMB [28] spectrum disfavors thermal relics of masses below 5 GeV. Therefore in order for χ^0 to have escaped these existing searches, its thermal annihilation cross-section has to be somewhat smaller than suggested by the dark-matter relic abundance measurements. This could happen if χ^0 is of Majorana type, so that its thermal annihilation cross-section is velocity suppressed. Alternatively, we could assume that χ^0 represents only a fraction of the total relic abundance. A third possibility is to lower the χ^0 mass below the sensitivity range of the present direct detection experiments, ignoring the cosmological bounds. If χ^0 is a Majorana particle, then only the axial current ($\bar{\chi}^0 \gamma_\mu \gamma_5 \chi^0$) part of (13) is allowed and current dark-matter searches (both direct and indirect) do not exclude a particle of this type with mass of a couple GeV or less (see e.g. [26, 27, 29]). In this case the correct dark-matter abundance can in principle still be obtained if χ^0 is produced non-thermally as in the asymmetric dark-matter scenarios [30], for example through the out-of-equilibrium decays of heavier particles, not included in our low-energy model.

V. CONCLUSIONS

Our knowledge of top-quark physics is still rather limited. As we have shown, it could well be that the $t\bar{t}$ sam-

ple analysed at the Tevatron is enriched by non-standard contributions coming from decays of a scalar top-partner, with the electric and color charge of the top quark. The top partners should be produced in pairs ($t\bar{t}$), have a mass slightly above m_t , and a large $\tilde{t} \rightarrow t\chi^0$ branching ratio, where χ^0 is a SM singlet with mass of a few GeV or less (escaping detection). With a proper tuning of the $q\tilde{t}\chi^0$ effective couplings, the additional $t\bar{t}$ pairs produced in this way could account for the large $A_{FB}^{t\bar{t}}$ observed at CDF.

The simple model we have proposed, where \tilde{t} and χ^0 are the only relevant new light states, is fully consistent with present high-energy data and contains a dark-matter candidate. The model requires a non-trivial flavour structure,³ such that \tilde{t} has large couplings to both $\chi^0 t_R$ and $\chi^0 u_R$, and a vanishing coupling to $\chi^0 c_R$. However, it is quite natural as far as the field content is concerned: a light scalar partner for the top is a natural candidate for a stabilization of the leading quadratic divergence on the SM Higgs mass term. In this paper we have pursued a bottom-up phenomenological approach and we have not analysed how this simple framework could be extended at high energies to be considered a viable extension of the SM; however, at this stage we see no obstacles to consider it as the low-energy side of a more ambitious supersymmetric model.

Interestingly, this framework can soon be tested in more detail at the LHC. In particular: i) the large E_T^{miss} of the sub-leading $\tilde{t} \rightarrow u\chi^0$ decay mode, and ii) the rise of the $pp \rightarrow t\bar{t} \rightarrow t\bar{t}$ QCD cross-section at large $m_{t\bar{t}}$, offer powerful tools to disproof or find evidences of this non-standard framework.

VI. ACKNOWLEDGMENTS

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VII. ADDED NOTE

After this work was completed, we become aware of a new anomaly reported by the CDF experiment, namely an excess in the production of jet pairs in association

³ A flavour-violating structure of this type has also been postulated in [31], although in a different framework, and could be attributed to a specific flavour symmetry.

with a W boson [32]. In particular, CDF reports an excess in the di-jet invariant mass (M_{jj}) distribution for $120 \lesssim M_{jj} \lesssim 160$ GeV. We note that our model predict an excess over the SM prediction in the $W + jj$ channel due to the sub-leading $\tilde{t} \rightarrow u + \chi$ decay mode: $\tilde{t}\tilde{t}^\dagger \rightarrow u\bar{t}$ ($t\bar{u}$) + $E_{\text{miss}} \rightarrow u\bar{b}$ ($b\bar{u}$) + $W^-(W^+) + E_{\text{miss}}$. We have simulated this decay chain and found that the resulting M_{jj} distribution is indeed peaked around 140 – 150 GeV for our preferred range of parameters. However, the size of the cross-section is substantially smaller (about 1/3) with respect to the central value

of the non-standard effect reported by CDF (after taking into account the experimental cuts applied in [32]). Moreover, the M_{jj} distribution predicted in our framework is substantially broader, with $\sigma(M_{jj}) \sim 100$ GeV with respect to the one reported in [32]. From these considerations we conclude that the inclusion of the $\tilde{t}\tilde{t}^\dagger$ production could partially improve the agreement with data also in the $W + jj$ channel, although a quantitative evaluation of this improvement would require a more accurate evaluation of detector efficiencies and resolutions, especially on the jet variables, that is beyond the scope of the present work.

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